SINP/TNP/00-21
hep-ph/0007242

Large CP violation in radiative B decays in supersymmetry without R-parity

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Abstract

We demonstrate that the R-parity breaking interactions within their current experimental upper bounds can give rise to large mixing-induced CP asymmetry in exclusive radiative B decays that may be detectable in the upcoming B factories.

PACS number(s): 11.30.Er, 13.25.Hw, 12.60.Jv, 11.30.Fs
Keywords: CP violation, B decays, R-parity violation

Supersymmetry is widely considered to be a leading candidate of physics beyond the standard model that may be realized in Nature. So much so that the search for its signals constitutes one of the main physics goals of the current and future colliders. On the other hand, there is an emerging consensus that new physics might show up through detectably large CP violation in B decays [1]. In this sense, it is reasonable to ask whether it can provide an indirect but an unmistakable signature of supersymmetry.

Since B decay modes are often masked by theoretical as well as experimental uncertainties, one way to establish the existence of new physics is to identify some characteristic signatures of it which can never be reproduced within the acceptable range of the standard model (SM) parameters. CP asymmetry in exclusive radiative B decays offers one such sensitive probe of physics beyond the SM. More explicitly, consider $B_q \to M^0\gamma$, where q is either d or s quark, and $M = \rho^0, \omega, \phi, K^{*0}$ (decaying as $K^{*0} \to K_s\pi^0$) is a self-conjugate meson with CP eigenvalue $\xi = \pm 1$. The CP asymmetry would arise due to the interference between mixing and decay. If ϕ_L denotes the weak phase associated with the $b_R \to s_L\gamma_L$ (i.e. $B_q \to M\gamma_L$) decay with amplitude A_L, while ϕ_R and A_R are corresponding quantities for $b_L \to s_R\gamma_R$ (i.e. $B_q \to M\gamma_R$), then the time-dependent mixing-induced CP asymmetry will be given by,

$$a_{\rm CP}(t) \equiv \frac{\Gamma(t) - \bar{\Gamma}(t)}{\Gamma(t) + \bar{\Gamma}(t)} = \xi A_M \sin(\phi_M - \phi_L - \phi_R)\sin(\Delta m t), \tag{1}$$

where, $A_M = 2|A_L A_R|/(|A_L|^2 + |A_R|^2)$, and ϕ_M is the phase of B_q-$\bar{B}_q$ mixing. In the SM, this asymmetry is small because of the following reason. At the quark level the interference occurs between $b_L \to q_R\gamma_R$ whose amplitude is proportional to m_q and the hermitian conjugate (through B_q-$\bar{B}_q$ mixing) of $b_R \to q_L\gamma_L$ whose amplitude is proportional to m_b, leading to $A_M \propto m_q/m_b$. As a result, this asymmetry is $\sim 1\%$ for $b \to d\gamma$ and $\sim 10\%$ for $b \to s\gamma$. The reason for this suppression is clearly the appearance of light quark mass m_q as a pre-factor with the $\bar{b}_L\sigma_{\mu\nu}q_R F^{\mu\nu}$ part (i.e. $b_L \to q_R\gamma_R$ decay) of the effective Hamiltonian *vis-a-vis* the appearance of m_b with the $\bar{b}_R\sigma_{\mu\nu}q_L F^{\mu\nu}$ part (i.e. $b_R \to q_L\gamma_R$ decay). Is it possible to avoid this suppression by going beyond the SM? In the left-right symmetric model this asymmetry can go up to 50% [2] and in supersymmetric model to about 80% with large sfermion mixings [3]. In this paper, we try to answer this question in supersymmetric models *without* R-parity.

In the minimal supersymmetric standard model (MSSM), gauge invariance ensures neither the conservation of lepton number (L) nor that of baryon number (B). Defining R-parity in terms of L and B

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as $R = (-1)^{(3B+L+2S)}$, where S is the spin of the particle, one should in a general supersymmetric model allow R-parity-violating (RPV) couplings [4]. R is $+1$ for all SM particles and -1 for their superpartners. Even though any concrete evidence for the existence of RPV terms is still lacking, the recent observation of neutrino masses and mixings in solar and atmospheric neutrino data suggests that it would be premature to abandon the L-violating interactions. However, to avoid rapid proton decay one cannot simultaneously admit both L- and B-violating interactions and for this reason we impose B conservation by hand. The L-violating superpotential can be written as (with i, j, k as generation indices)

$$W_{\rm RPV} \equiv \frac{1}{2}\lambda_{ijk} L_i L_j E^c_k + \lambda'_{ijk} L_i Q_j D^c_k + \mu_j L_j H_u, \tag{2}$$

where $\lambda_{ijk} = -\lambda_{jik}$. Here L_i and Q_i are lepton and quark doublet superfields, E^c_i and D^c_i are charged lepton and down quark singlet superfields, and H_u is the Higgs superfield that gives mass to up-type quarks. The trilinear λ_{ijk}-couplings and bilinear μ_i mass parameters are not relevant for our purpose and from now on we consider only the trilinear λ'_{ijk}-induced interactions.

Tight constraints on the sizes of these couplings have been placed from the consideration of neutrinoless double beta decay, ν_e-Majorana mass, charged-current universality, $e-\mu-\tau$ universality, ν_μ deep-inelastic scattering, atomic parity violation, τ decays, D and K decays, Z decays, etc. Product of two couplings at a time have been constrained from $\mu - e$ conversion, $\mu \to e\gamma$, $b \to s\gamma$, B decays into two charged leptons, $K_L - K_S$ and $B_q - \overline{B_q}$ $(q = d, s)$ mass differences, etc. For a collection of all these limits see [5].

Let us now turn our attention to the $b \to s\gamma$ amplitude in RPV models (Needless to add that a similar analysis can be carried out for $b \to d\gamma$) [6]. The aim is to generate unsuppressed diagrams for $b_L \to s_R\gamma_R$. The RPV diagrams are shown in Fig. 1. The trilinear L-violating couplings involved are λ'_{ij2} and λ'_{ij3}.

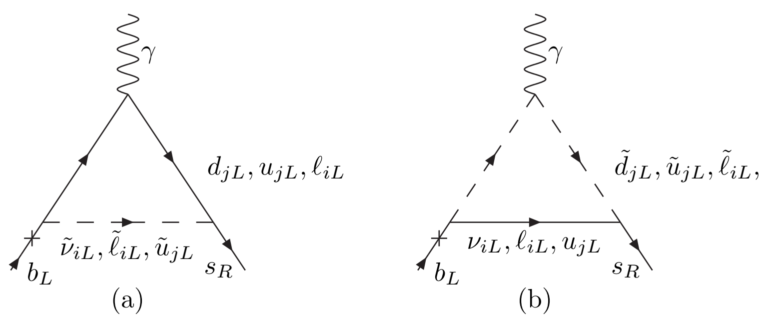

Fig. 1: The diagrams for $b_L \to s_R\gamma_R$ induced by λ'_{ij2} and λ'_{ij3} couplings.

(i) Figure 1a is a generic diagram in which the photon couples to the internal fermion line. It follows from the RPV superpotential, and as shown in the Figure 1a, that the fermion - scalar combinations in this case are $(d_{jL}, \tilde\nu_{iL})$, $(u_{jL}, \tilde\ell^-_{iL})$, and $(\ell^-_{iL}, \tilde u_{jL})$.
(ii) Figure 1b is a similar to diagram where the photon is attached to the internal scalar. This time the fermion - scalar combinations are $(\nu_{iL}, \tilde d_{jL})$, $(\ell^-_{iL}, \tilde u_{jL})$, and $(u_{jL}, \tilde\ell^-_{iL})$.

In Fig. 1, although chiral flips are marked only on the external b quark lines thus leading to a contribution proportional to m_b, similar diagrams can be drawn with chiral flips on the external s quark lines as well (leading to $b_R \to s_L\gamma_L$), the latter contribution being proportional to m_s. The sum of amplitudes, taking into account all fermion-sfermion combinations in Fig. 1, at the electroweak scale, is given by

$$i\Gamma^{(1)}_\mu = -\sum_{ij} \frac{ie}{32\pi^2} \frac{\lambda'_{ij2}\lambda'^*_{ij3}}{m_W^2} \bar s(m_b P_L + m_s P_R) i\sigma_{\mu\nu} q^\nu b F^{ij}_R \ , \tag{3}$$

with $F^{ij}_R = Q_d F^{ijd} + Q_u F^{iju} + Q_{\ell^-} F^{ij\ell^-}$, where

$$F^{ijd} = \frac{m_W^2}{m^2_{\tilde\nu_i}} G_1\left(\frac{m^2_{d_j}}{m^2_{\tilde\nu_i}}\right) + \frac{m_W^2}{m^2_{\tilde d_j}} G_2\left(\frac{m^2_{\nu_i}}{m^2_{\tilde d_j}}\right) \ , \tag{4}$$

$$F^{iju} = \frac{m_W^2}{m_{\tilde{\ell}_i^-}^2} G_1\left(\frac{m_{u_j}^2}{m_{\tilde{\ell}_i^-}^2}\right) + \frac{m_W^2}{m_{\tilde{u}_j}^2} G_2\left(\frac{m_{\ell_i^-}^2}{m_{\tilde{u}_j}^2}\right) , \tag{5}$$

$$F^{ij\ell^-} = \frac{m_W^2}{m_{\tilde{u}_j}^2} G_1\left(\frac{m_{\ell_i^-}^2}{m_{\tilde{u}_j}^2}\right) + \frac{m_W^2}{m_{\tilde{\ell}_i^-}^2} G_2\left(\frac{m_{u_j}^2}{m_{\tilde{\ell}_i^-}^2}\right) . \tag{6}$$

Here the Inami-Lim functions [7] are

$$G_1(x) = \xi_1(x) - \xi_2(x) \;, G_2(x) = -\xi_0(x) + 2\xi_1(x) - \xi_2(x) \;, \tag{7}$$

$$\xi_n(x) \equiv \int_0^1 \frac{z^{n+1}\mathrm{d}z}{1+(x-1)z} = -\frac{\ln x + (1-x) + \cdots + \frac{(1-x)^{n+1}}{n+1}}{(1-x)^{n+2}} . \tag{8}$$

The charge Q factor in the expression for F_R^{ij} reflects the charge of the internal line where the photon leg is attached to. The function G_1 (G_2) corresponds to the case that the photon attaches to the internal fermion (sfermion). The argument $m_j^2/m_{\tilde{i}}^2$ in functions G's indicates internal lines of the fermion of generation j and the sfermion of generation i. Explicitly,

$$G_1(x) = \frac{2+5x-x^2}{6(1-x)^3} + \frac{x\ln x}{(1-x)^4} \;, \quad G_2(x) = -\frac{1-5x-2x^2}{6(1-x)^3} + \frac{x^2\ln x}{(1-x)^4} \;. \tag{9}$$

Except for $j=3$ in the combination $\lambda'_{ij2}\lambda'^*_{ij3}$, in which case there could be an internal top quark line, we can set the argument $x \to 0$ to obtain

$$G_1(0) = \frac{1}{3} \;, \quad G_2(0) = -\frac{1}{6} \;. \tag{10}$$

In the limit of a common mass $\tilde{m}$ for all sfermions, we then have

$$F_R^{ij} \simeq -\frac{m_W^2}{9\tilde{m}^2} \;. \tag{11}$$

To be precise, we should treat the heavy top sector separately according to the detailed formulas. However, we use Eq. (11) only for simplicity.

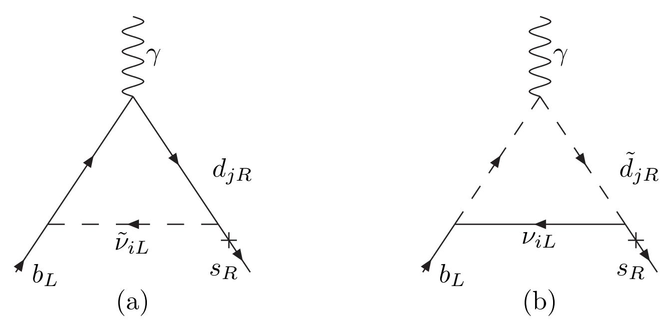

Fig. 2: The diagrams for $b_L \to s_R\gamma_R$ induced by λ'_{i3j} and λ'_{i2j} couplings.

In Fig. 2, we show graphs of another type, with couplings λ'_{i3j} and λ'_{i2j}. This time we do not explicitly exhibit the diagrams where the chiral flips are on the b quark lines (corresponding to $b_R \to s_L\gamma_L$). They give rise to the amplitude

$$i\Gamma_\mu^{(2)} = -\sum_{ij} \frac{ie}{32\pi^2} \frac{\lambda'_{i3j}\lambda'^*_{i2j}}{m_W^2} \bar{s}(m_b P_R + m_s P_L) i\sigma_{\mu\nu} q^\nu b F_L^{ij} \;, \tag{12}$$

Assuming degenerate masses for the left-handed and right-handed squarks, we obtain $F_L^{ij} = Q_d F^{ijd}$.

Including the two types of graphs together with the SM contribution, we set up the effective Hamiltonian responsible to the process $b \to s\gamma$,

$$\mathcal{H}^{\rm eff} = \frac{1}{64\pi^2 m_W^2} \sum_h \bar{s}\sigma_{\alpha\beta} \left(c_{7h} e F^{\alpha\beta} + c_{8h} g_s T^a G^{a\alpha\beta} \right) (m_b P_{-h} + m_s P_h) b \ , \tag{13}$$

with the chirality index $h = \pm$ (or R, L) indicating $P_\pm = \frac{1}{2}(1 \pm \gamma_5)$. The term involving c_{8h} corresponds to the gluonic dipole contribution. At the electroweak scale m_W, the Wilson coefficients are given by the short-distance amplitudes outlined above,

$$c_{7,8R}(m_W) = \sum_{ij} \lambda'_{ij2} \lambda'^*_{ij3} F^{ij}_{7,8R} \ . \tag{14}$$

It is straightforward to see $F^{ij}_{7R} = F^{ij}_R$ and $F^{ij}_{8R} = F^{ijd} + F^{iju}$. Also,

$$c_{7,8L}(m_W) = c^{\rm SM}_{7,8L} + \sum_{ij} \lambda'_{i3j} \lambda'^*_{i2j} F^{ij}_{7,8L}, \quad \text{with} \quad F^{ij}_{7L} = Q_d F^{ijd}, \ F^{ij}_{8L} = F^{ijd}. \tag{15}$$

The SM contributions are well known to be [7] (with $x_t = m_t^2/m_W^2$)

$$c^{\rm SM}_{7L}(m_W) = g^2 V_{tb} V^*_{ts} x_t \left[\frac{7 - 5x_t - 8x_t^2}{12(1-x_t)^3} + \frac{2x_t - 3x_t^2}{2(1-x_t)^4} \ln x_t \right] \ , \tag{16}$$

$$c^{\rm SM}_{8L}(m_W) = g^2 V_{tb} V^*_{ts} x_t \left[\frac{2 + 5x_t - x_t^2}{4(1-x_t)^3} + \frac{3x_t \ln x_t}{2(1-x_t)^4} \right] . \tag{17}$$

For our estimation, we evaluate the Wilson coefficients at the m_b scale by the simplified renormalization group evolution [8],

$$c_{7h}(m_b) = \eta^{\frac{16}{23}} c_{7h}(m_W) + \tfrac{8}{3} \left(\eta^{\frac{14}{23}} - \eta^{\frac{16}{23}} \right) c_{8h}(m_W) + N^{\rm SM}_h \ , \qquad \text{for } h = L \text{ or } R, \tag{18}$$

with $\eta = \alpha_s(m_W)/\alpha_s(m_b) = 0.56$. The leading log QCD corrections in the SM are given by,

$$N^{\rm SM}_L = g^2 V_{tb} V^*_{ts} \tfrac{464}{513} \left(\eta^{-\frac{3}{23}} - \eta^{\frac{16}{23}} \right) \ , \quad N^{\rm SM}_R = 0 \ . \tag{19}$$

For the purpose of *order-of-magnitude* estimate, we have ignored other leading logarithmic contributions arising from additional operators of RPV origin, which have been outlined in Refs. [9, 10].

Numerically, for the a common mass $\tilde{m}$, same for squarks and sleptons, we have

$$c_{7R}(m_b) = D \left(\frac{m_W^2}{\tilde{m}^2} \right) \sum_{ij} \lambda'_{ij2} \lambda'^*_{ij3} \ , \tag{20}$$

$$c_{7L}(m_b) = g^2 V_{tb} V^*_{ts} E + \frac{1}{2} D \left(\frac{m_W^2}{\tilde{m}^2} \right) \sum_{ij} \lambda'_{i3j} \lambda'^*_{i2j} \ , \tag{21}$$

for $E \simeq 0.65$ and

$$D = \tfrac{8}{9} \eta^{\frac{14}{23}} - \eta^{\frac{16}{23}} \simeq -0.044 \ . \tag{22}$$

The size of the mixing-induced CP asymmetry will be controlled by the relative magnitude of A_R and A_L (for the exact calculation, one should use Eq. (1)). This is given by the ratio

$$R \equiv \left| \frac{A_R}{A_L} \right| = \left| \frac{c_{7R} + (m_s/m_b) c_{7L}}{c_{7L} + (m_s/m_b) c_{7R}} \right| \ , \tag{23}$$

with the Wilson coefficients evaluated at the m_b scale.

$\lambda'_{ij2}\lambda'_{ij3}$	Upper limits	$\lambda'_{i2j}\lambda'_{i3j}$	Upper limits
(112)(113)	4.0×10^{-4}	(121)(131)	4.0×10^{-4}
(212)(213)	2.5×10^{-3}	(221)(231)	1.1×10^{-2}
(312)(313)	2.5×10^{-3}	(321)(331)	6.1×10^{-2}
(122)(123)	8.0×10^{-4}	(122)(132)	$3.3 \times 10^{-3}(\dagger)$
(222)(223)	$2.7 \times 10^{-3}(*)$	(222)(232)	$3.3 \times 10^{-3}(\dagger)$
(322)(323)	$2.7 \times 10^{-3}(*)$	(322)(332)	$3.3 \times 10^{-3}(\dagger)$
(132)(133)	2.8×10^{-4}	(123)(133)	2.8×10^{-5}
(232)(233)	$2.5 \times 10^{-3}(*)$	(223)(233)	$3.3 \times 10^{-3}(\dagger)$
(332)(333)	$2.5 \times 10^{-3}(*)$	(323)(333)	$3.3 \times 10^{-3}(\dagger)$

Table 1: The existing 1-σ upper limits on the λ' product couplings that enter into the expressions of c_{7R} and c_{7L}. The meaning of $*$ and $\dagger$ symbols have been described in the text. A common 100 GeV mass for all scalars have been assumed while obtaining those limits.

The existing upper limits (1-σ) on the magnitudes of $\lambda'_{ij2}\lambda'_{ij3}$ and $\lambda'_{i2j}\lambda'_{i3j}$ combinations have been listed in Table 1 (while deriving limits the couplings have been assumed to be real in most cases). We have assumed a common mass of 100 GeV for whichever scalar is exchanged. In most cases the best bounds have been obtained by multiplying the individual upper limits on the respective couplings, listed in Ref. [11]. Note, B_s–$\bar{B}_s$ mixing (actually taking its lower limit in a conservative sense) constrains the combinations $\lambda'_{ij2}\lambda'_{ij3}$ via sneutrino mediated one-loop box graphs [12]. In some cases, depending on the generation indices of the internal quarks and squarks, these bounds are stronger than those obtained by multiplying individual bounds. These have been marked $(*)$ in Table 1. The bounds marked $(\dagger)$ have been derived [10] from $B \to \phi K$ decays and semi-leptonic B decays (Note, the authors of Ref. [10] have missed the color suppression factor in the RPV amplitudes. We have included this factor when quoting bounds in Table 1).

As far as CP asymmetry is concerned, the size of $\sum_{ij} \lambda'^{*}_{i2j}\lambda'_{i3j}$ is not so significant. The reason is that this product (even set at its upper limit) multiplied by $F_{7,8L}$ adds to a much larger contribution from $c^{\rm SM}_{7,8L}$ (see Eq. (15)). The size of the other combination, namely $\sum_{ij} \lambda'_{ij2}\lambda'^{*}_{ij3}$, is however crucial for our prediction for CP asymmetry. It follows from Table 1 that we can, as a rough estimate, take the magnitude of this combination to be $\sim \pm 0.01$. This yields rather small $R \sim 3\%$. Such a tiny effect can hardly be interpreted as a signal of new physics. On the other hand, if one assumes that the squarks are much heavier than the sleptons, as happens in the gauge-mediated supersymmetry breaking models [13], then in the RPV amplitudes one can consider only the slepton-mediated graphs, as the squark-mediated ones sufficiently decouple to leave no numerical impact. *For our remaining discussions, we stick to the latter scenario.* In this case the factor of $\frac{1}{2}$ in Eq. (21) becomes 1 and the variable D in Eq. (20) is replaced by

$$D' = \tfrac{16}{9}\eta^{\frac{14}{23}} - \tfrac{3}{2}\eta^{\frac{16}{23}} \simeq 0.25 \ , \tag{24}$$

which is rather large compared to D. This time, with $|\sum_{ij} \lambda'_{ij2}\lambda'^{*}_{ij3}| \sim 0.01$, we obtain R as large as 16%. Note that if we set this combination to be 0.04 (0.05), which is not very unrealistic, as the limits have been set only on an order-of-magnitude basis, R may go up to $\sim 70\%$ (90%), which is a rather large effect.

Although in the case of CP asymmetry the impact of $\sum_{ij} \lambda'^{*}_{i2j}\lambda'_{i3j}$ is insignificant as noted earlier, this combination, both in magnitude and phase, may turn out to be quite significant in the total branching ratio (BR) for $b \to s\gamma$. Since the BR is proportional to $(c^2_{7R} + c^2_{7L})$, the RPV contribution to c_{7L} in Eq. (21) interferes with a rather large SM contribution and hence appears linearly, while the RPV contribution to c_{7R} (being the only contribution) appears only quadratically. Also, within c_{7L}, the interference between the SM and the RPV pieces may be either constructive or destructive depending on their relative phase. The possibility of this partial cancellation may enable us to admit larger values of RPV couplings required to generate a sizable CP asymmetry, while stay perfectly consistent with the BR constraints. In Fig. 3, as an illustrative example. we exhibit our prediction for the BR as a function of $|\lambda^2|$, where $\lambda^2 \equiv \sum_{ij} \lambda'_{ij2}\lambda'^{*}_{ij3} =$

$\sum_{ij} \lambda'_{i3j}\lambda'^{*}_{i2j}$. The different angles refer to the relative phase between the two pieces in Eq. (21). In Fig. 4, we demonstrate how large CP asymmetry (only its absolute value) can be generated as a function of $|\lambda^2|$ without violating the BR constraint. For the latter, we used Eq. (1) with the sine of the phase combination set to unity. Here we wish to make the following remark. While drawing the new physics curves in Fig. 3, we took a conservative approach that the SM reference line is at its present mean value. One might as well place the SM reference line close to its lower limit so that larger values of $|\lambda^2|$ could be allowed leading to a possibility of larger CP asymmetry.

Before closing, we highlight the following salient features:

1. The sign of the RPV induced CP asymmetry is arbitrary, since the signs of the λ' couplings are *a priori* unknown. It is therefore possible to have a CP asymmetry not only large compared to the SM prediction but also with an opposite sign. The occurance of a sign flip, in particular, may constitute an unmistakable signature of new physics [14].

2. In principle, the λ' couplings need not be complex in order to generate CP asymmetry. More precisely, in Eq. (1) even with $\phi_R = 0$ (i.e. with real λ' couplings), one may obtain a non-zero sine function from the B-$\bar{B}$ mixing phase ϕ_M or the SM decay phase ϕ_L. Nevertheless, a non-zero ϕ_R may trigger a much larger CP violating effect.

3. Here we present a simple way of understanding how large CP asymmetry can one predict without violating the branching ratio constraint. The present experimental BR on $B \to X_s\gamma$ is $(3.15 \pm 0.54) \times 10^{-4}$ [15], while the present SM next-to-leading order estimate for BR $(b \to s\gamma)$ is $(3.28 \pm 0.33) \times 10^{-4}$ [16]. To check whether our prediction of a large CP asymmetry at all contradicts the constraints on the BR, let us consider the following illustrative example. Assume that new physics contributes only to c_{7R}. In our case, this amounts to saying that $\sum_{ij} \lambda'_{ij2}\lambda'^{*}_{ij3}$ is non-zero, while $\sum_{ij} \lambda'_{i3j}\lambda'^{*}_{i2j}$ is vanishingly small. Recall, to a good approximation, $c_{7R}/c_{7L} \equiv \alpha$ is a quantity that controls the size of the CP asymmetry (for an exact expression, see Eq. (1)). Then the BR, being proportional to $(c_{7R}^2 + c_{7L}^2)$, is modified by an overall factor $(1 + \alpha^2)$. Now to demonstrate how large α one can tolerate, let us imagine a situation when the SM prediction of the BR is at its present 90% CL lower limit, i.e. 2.74×10^{-4}, and the experimental BR settles at its 90% CL upper limit, i.e. 4.04×10^{-4}. This means that within the current experimental and theoretical uncertainties a 50% ($= \alpha^2$) new physics induced enhancement of BR is perfectly tolerable, and hence a CP asymmetry as large as $2\alpha/(1+\alpha^2) \sim 100\%$ (putting $\alpha = 0.7$ in Eq. (1)) is very much consistent with the present constraints on the BR at 90% CL.

4. The RPV models are better placed than the left-right symmetric model in generating large CP asymmetry. The reason is that the R-parity conserving MSSM contribution to c_{7L} may interfere destructively with the SM amplitude in some region of the parameter space [17]. Just like the RPV contribution to c_{7L}, this also helps to keep the BR under control while admitting large RPV couplings responsible for large CP asymmetry. More specifically, the charged Higgs (chargino) loops interfere constructively (destructively) with the SM W loops, and we know that in the exact supersymmetric limit the net amplitude vanishes [18].

5. The combinations $\lambda'_{i12}\lambda'^{*}_{i13}$, within their current experimental bounds, not only induce large CP asymmetry in $b \to s\gamma$ decay, but also enhance the direct CP asymmetry in $B^{\pm} \to \pi^{\pm}K$ channels [19]. Hence if upcoming B factories indicate large CP asymmetry in both these (uncorrelated) channels, and perhaps with signs opposite to their respective SM predictions, it may constitute an unambiguous signal of the RPV scenario. We must admit though that the extraction of γ from $B \to \pi K$ would be even more difficult in the presence of new physics contamination.

GB thanks the following centers for hospitality where portions of the work were done: NCTS, Hsinchu, Taiwan; CERN Theory Division, Geneva, Switzerland; and MPI Kernphysik, Heidelberg, Germany. DC wishes to thank SLAC Theory Group and U.S. DOE for hospitality while revising the final manuscript.

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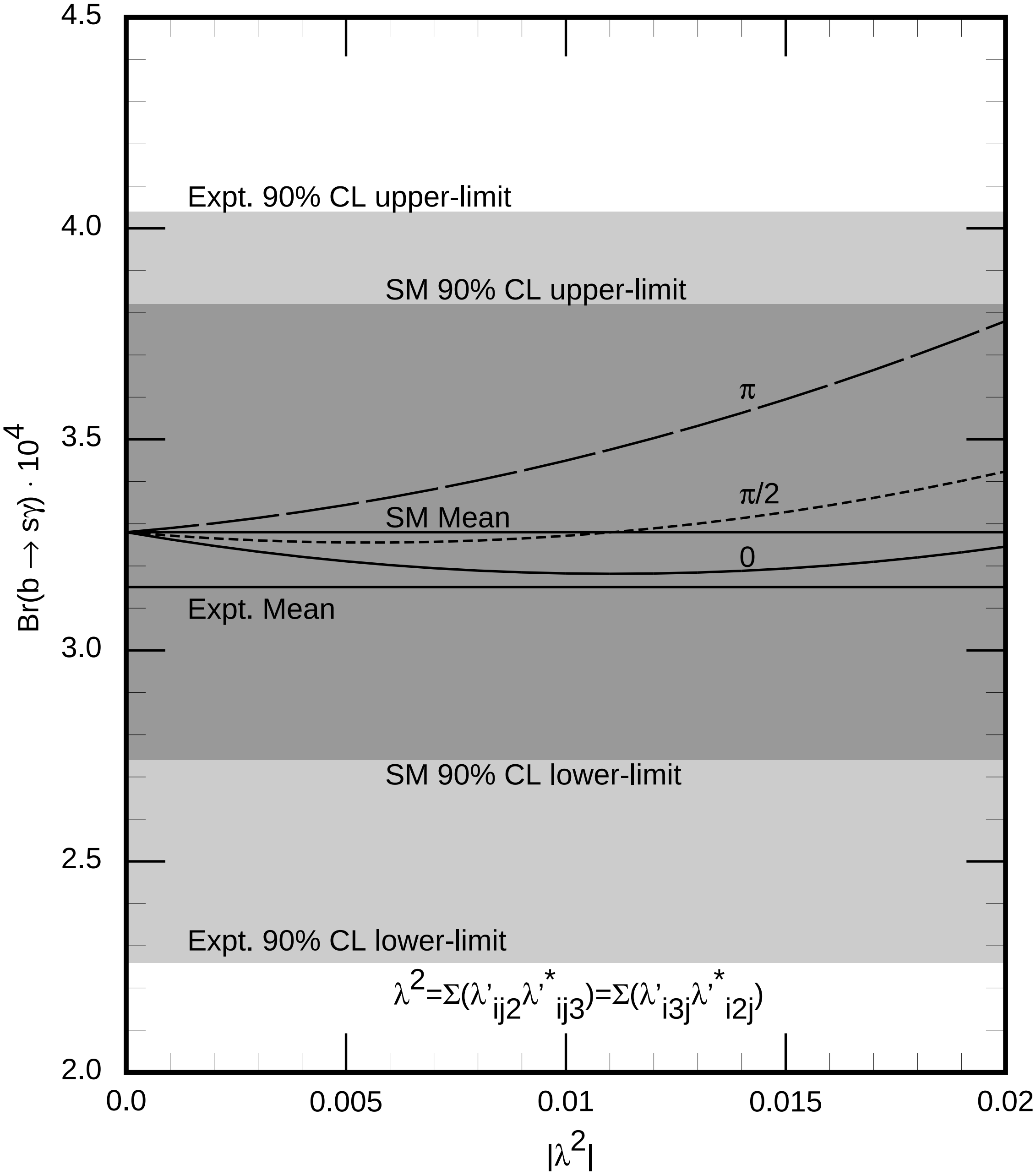

Fig. 3: Br($b \to s\gamma$) versus $|\lambda^2|$ for various choices of the relative phase between $\sum_{ij} \lambda'_{i3j}\lambda'^{*}_{i2j}$ and $V_{tb}V^{*}_{ts}$. The SM prediction and the experimentally allowed region are shaded.

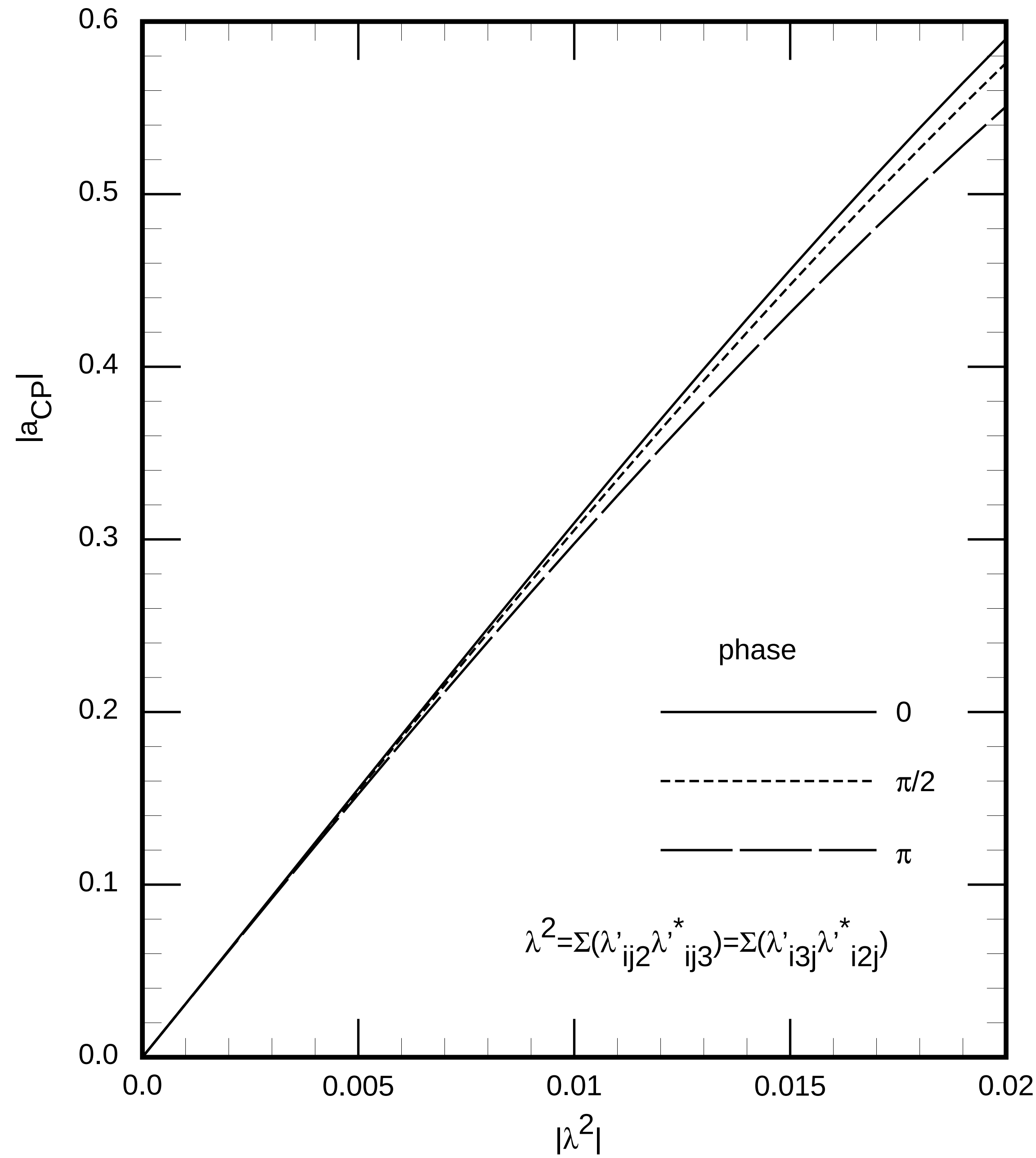

Fig. 4: The magnitude of the CP asymmetry for various choices of the relative phase between $\sum_{ij} \lambda'_{i3j}\lambda'^{*}_{i2j}$ and $V_{tb}V_{ts}^{*}$. Here we assume that both the sine factors are unity in Eq. (1).